NEUTRINO MASSES AND MIXING IN SPLIT SUPERSYMMETRY

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We investigate the possibility to describe the neutrino masses and mixing angles in the context of Split Supersymmetric scenarios. All relevant contributions coming from the R-parity violating terms to the neutrino mass matrix up to one-loop level are computed, showing the importance of the Higgs boson one-loop correction, which has been overlooked in previous studies. We conclude that it is possible to explain all neutrino masses and mixings in Split Supersymmetry with bilinear R-Parity violating interactions.

I. INTRODUCTION

Supersymmetric extensions of the Standard Model (SM) have been considered as one of the most serious candidates for the physics beyond the Standard Model. In the last years different supersymmetric scenarios have been studied extensively. We mention low-energy SUSY [1], where the supersymmetric scale is around TeV, and Split SUSY where all the scalars, except for one Higgs doublet, are very heavy [2].

In both supersymmetric scenarios mentioned above it is possible to achieve unification of the gauge interactions at the high scale and the lightest supersymmetric particle (LSP) could be a natural candidate to describe the Cold Dark Matter in the Universe once the so-called R-parity is imposed as an exact symmetry of the theory. However, in SPLIT SUSY scenarios, ignoring the hierarchy problem, most of the unpleasant aspects of low-energy SUSY, such as excessive flavour and CP violation, and very fast dimension 5 proton decay, are eliminated.

It is very-well known that in general interactions which break the lepton or baryon number (or R-parity) are present in any SUSY extension of the SM. Therefore, we have the possibility to describe the neutrino masses and mixing [3], and we have to understand the predictions for proton stability [4] in this context.

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For several phenomenological aspects of R-parity violating interactions see [5].

The possibility to describe the neutrino properties with R-parity violating interactions in the context of the Minimal Supersymmetric Standard Model has been studied in detail for several groups in the context of low energy supersymmetry (See for example [6, 7]). Recently, in the context of SPLIT SUSY the possibility to describe the masses and mixing of neutrinos has been studied in reference [8], where the authors concluded that it is not possible to use the R-parity bilinear terms alone to describe the neutrino properties. See also reference [9] for a nice model that uses a radiative see-saw mechanism [10] in the context of split SUSY.

In this work we re-examine the possibility to describe the properties of neutrinos using the R-parity violating interactions in the context of split supersymmetric scenarios. We show that it is possible to describe the neutrino properties using all relevant interactions once the heavy scalars are integrated out. Computing all contributions up to one-loop level, we find an example solution where it is shown that all constraints from neutrino experiments on the R-parity violating interactions are satisfied.

We conclude that it is sufficient to use the bilinear terms alone in order to explain the neutrino masses and mixing angles, and that trilinear R-Parity violating (TRpV) couplings are essentially irrelevant due to the large mass of the scalars. The key element, overlooked in previous studies, is the Higgs boson loop together with gauginos, which can couple to neutrinos via the mixing with higgsinos.

II. R-PARITY VIOLATION AND NEUTRINO MASSES IN SPLIT SUSY

As we know in any supersymmetric extension of the Standard Model there are interactions terms which break the so-called R-parity. The R-parity is defined as $R = (-1)^{3(B-L)+2S}$, where L, B, and S are the lepton and baryon number, and the spin, respectively. Usually this symmetry is considered as an exact symmetry of the minimal supersymmetric extension of the Standard Model in order to avoid problems with proton decay and at the same time there is a possibility to have the lightest supersymmetric particle as a good candidate for the Cold Dark Matter of the Universe. In the context of SPLIT SUSY these issues have been studied in [11]. See reference [12] for the possibility to having R-parity as an exact symmetry coming from grand unified theories.

In this work we focus on a particular supersymmetric scenario, Split SUSY. Integrating out the heavy scalars all possible R-parity conserving interactions in split supersymmetric scenarios are given by [2]:

$$\begin{aligned}\mathcal{L}^{Split}_{SUSY} \;=\; & \mathcal{L}^{Split}_{Kinetic} \;+\; m^2 H^\dagger H - \frac{\lambda}{2}(H^\dagger H)^2 - \Big[Y_u \overline{q}_L u_R i\sigma_2 H^* \;+\; Y_d \overline{q}_L d_R H \;+\; Y_e \overline{l}_L e_R H \\ & + \frac{M_3}{2}\widetilde{G}\widetilde{G} \;+\; \frac{M_2}{2}\widetilde{W}\widetilde{W} \;+\; \frac{M_1}{2}\widetilde{B}\widetilde{B} \;+\; \mu \widetilde{H}_u^T i\sigma_2 \widetilde{H}_d \\ & + \tfrac{1}{\sqrt{2}} H^\dagger (\tilde{g}_u \sigma \widetilde{W} \;+\; \tilde{g}'_u \widetilde{B})\widetilde{H}_u \;+\; \tfrac{1}{\sqrt{2}} H^T i\sigma_2 (-\tilde{g}_d \sigma \widetilde{W} + \tilde{g}'_d \widetilde{B})\widetilde{H}_d + \text{h.c.} \Big]\end{aligned} \tag{1}$$

where H is the SM Higgs-doublet. In the above equations we have the SM fields q_L, u_R, d_R, l_L, e_R and the superpartners of the Higgs and gauge bosons in the MSSM. Following our notation $\widetilde{G}$, $\widetilde{W}$, and $\widetilde{B}$ are the gauginos associated to the $SU(3)$, $SU(2)$, and $U(1)$ gauge groups, respectively. While $\widetilde{H}_u$ and $\widetilde{H}_d$ correspond to the up and down higgsinos. The parameters in eq. (1) are the following: m is the Higgs mass parameter, λ is the Higgs quartic self coupling; Y_u, Y_d, and Y_e are the Yukawa couplings; M_3, M_2, and M_1 are the gaugino masses, μ the higgsino mass, and $\tilde{g}_u$, $\tilde{g}'_u$, $\tilde{g}_d$, and $\tilde{g}'_d$ are trilinear couplings between the Higgs boson, gauginos, and higgsinos.

The Higgs-gaugino-higgsino couplings in eq. (1) satisfy matching conditions at the scale $\tilde{m}$. Above this scale, the theory is supersymmetric and the squarks, sleptons, and heavy Higgs doublet have a mass assumed to be nearly degenerate and equal to $\tilde{m}$. The supersymmetric lagrangian includes the terms,

$$\mathcal{L}_{susy} \ni \mu \widetilde{H}_u^T i\sigma_2 \widetilde{H}_d - \frac{H_u^\dagger}{\sqrt{2}}\left(g\sigma\widetilde{W} + g'\widetilde{B}\right)\widetilde{H}_u - \frac{H_d^\dagger}{\sqrt{2}}\left(g\sigma\widetilde{W} - g'\widetilde{B}\right)\widetilde{H}_d \tag{2}$$

which implies the following boundary conditions at $\tilde{m}$:

$$\begin{aligned}\tilde{g}_u(\tilde{m}) = g(\tilde{m}) sin\beta(\tilde{m}) \quad , \qquad & \tilde{g}_d(\tilde{m}) = g(\tilde{m}) cos\beta(\tilde{m}) \\ \tilde{g}'_u(\tilde{m}) = g'(\tilde{m}) sin\beta(\tilde{m}) \quad , \qquad & \tilde{g}'_d(\tilde{m}) = g'(\tilde{m}) cos\beta(\tilde{m})\end{aligned} \tag{3}$$

where $g(\tilde{m})$ and $g'(\tilde{m})$ are the gauge coupling constants evaluated at the scale $\tilde{m}$. At the same time the angle β is the mixing angle between the two Higgs doublets H_d and H_u of the supersymmetric model. In terms of these two supersymmetric Higgs doublets of the MSSM, the light fine-tuned Higgs doublet H in the low energy effective model is $H = -i\sigma_2 H_d^* c_\beta(\tilde{m}) + H_u s_\beta(\tilde{m})$.

At low energy, the gauge couplings satisfy the RGE given in [2]. In particular, $\tan\beta(\tilde{m})$ evolves into the two independent ratios $\tilde{g}_u/\tilde{g}_d$ and $\tilde{g}'_u/\tilde{g}'_d$ which satisfy:

$$\begin{aligned}\frac{\tilde{g}_u}{\tilde{g}_d}(m_W) \;\approx\; & tan\beta(\tilde{m})\left\{1 + \frac{cos(2\beta)}{64\pi^2}(7g^2 - 3{g'}^2)\Big|_{\tilde{m}} \ln\frac{\tilde{m}}{m_W}\right\} \equiv tan\beta \\ \frac{\tilde{g}'_u}{\tilde{g}'_d}(m_W) \;\approx\; & tan\beta(\tilde{m})\left\{1 - \frac{cos(2\beta)}{64\pi^2}(9g^2 + 3{g'}^2)\Big|_{\tilde{m}} \ln\frac{\tilde{m}}{m_W}\right\} \equiv tan'\beta\end{aligned} \tag{4}$$

For simplicity we call them t_β and t'_β, respectively. Now, let us define the following quantities,

$$\tilde{g}^2 \equiv \tilde{g}_u^2(m_W) + \tilde{g}_d^2(m_W)\,, \qquad \tilde{g}'^2 \equiv \tilde{g}'^2_u(m_W) + \tilde{g}'^2_d(m_W) \tag{5}$$

the relation between the former quantities and $g^2(\tilde{m})$ and $g'^2(\tilde{m})$ can be easily found from the RGE in [2].

As we mentioned before in SPLIT SUSY scenarios at low energy we have the SM fields, the charginos and neutralinos. Using the above notation the chargino mass matrix is given by:

$$\mathbf{M}_{\chi^+} = \begin{bmatrix} M_2 & \frac{1}{\sqrt{2}} v\tilde{g}s_\beta \\ \frac{1}{\sqrt{2}} v\tilde{g}c_\beta & \mu \end{bmatrix} \tag{6}$$

while the neutralino mass matrix reads as:

$$\mathbf{M}_{\chi^0} = \begin{bmatrix} M_1 & 0 & -\frac{1}{2}\tilde{g}'c'_\beta v & \frac{1}{2}\tilde{g}'s'_\beta v \\ 0 & M_2 & \frac{1}{2}\tilde{g}c_\beta v & -\frac{1}{2}\tilde{g}s_\beta v \\ -\frac{1}{2}\tilde{g}'c'_\beta v & \frac{1}{2}\tilde{g}c_\beta v & 0 & -\mu \\ \frac{1}{2}\tilde{g}'s'_\beta v & -\frac{1}{2}\tilde{g}s_\beta v & -\mu & 0 \end{bmatrix} \tag{7}$$

Now, since we are interested in the possibility to describe the neutrino masses in Split-SUSY, we write all relevant R-Parity violating interactions:

$$\mathcal{L}_{RpV}^{Split} = \epsilon_i \widetilde{H}_u^T i\sigma_2 L_i \; - \; \frac{1}{\sqrt{2}} a_i H^T i\sigma_2 (-\tilde{g}_d \sigma \widetilde{W} + \tilde{g}'_d \widetilde{B}) L_i \; + \; h.c. \tag{8}$$

where ϵ_i are the parameters that mix higgsinos with leptons, and a_i are dimensionless parameters that mix gauginos with leptons. Notice that the first term is the usual bilinear term, while the last two terms are obtained once we integrate out the sleptons using the bilinear soft terms ($\tilde{L}_i H_u$) which break explicitly R-parity. As it is well-known we can also write the usual R-parity violating trilinear terms ($\hat{Q}\hat{D}^C\hat{L}$, $\hat{L}\hat{L}\hat{E}^C$). However, since the sfermions are very heavy in SPLIT SUSY and the contributions to the neutrino mass matrix coming from those terms are at one-loop level, those interactions cannot play any important role. The previous lagrangian is to be compared with the supersymmetric one valid above the scale $\tilde{m}$, from which we highlight the terms,

$$\mathcal{L}_{susy} \ni \epsilon_i \widetilde{H}_u^T i\sigma_2 L_i - \frac{\widetilde{L}_i^\dagger}{\sqrt{2}} \left(g\sigma^a \widetilde{W}^a - g'\widetilde{B} \right) L_i \tag{9}$$

in this case replacing $\widetilde{L}_i^*$ by $s_i i\sigma_2 H$ we get the matching condition:

$$a_i(\widetilde{m}) = \frac{s_i(\widetilde{m})}{\cos\beta(\widetilde{m})} \tag{10}$$

where the parameters $s_i(\widetilde{m})$ represent the amount of slepton $\widetilde{L}_i$ in the low energy Higgs H. Using eq. (8), after the Higgs acquires a vev, we find the relevant terms for neutrino masses:

$$\mathcal{L}_{RpV}^{Split} = \left[\epsilon_i \widetilde{H}_u^0 + \frac{1}{2} a_i v \left(\tilde{g} c_\beta \widetilde{W}_3 - \tilde{g}' c'_\beta \widetilde{B} \right) \right] \nu_i \; + \; h.c. \; + \; \ldots \tag{11}$$

where v is the vacuum expectation value of the SM-like Higgs field H. Knowing all R-parity violating interactions we can write the neutralino/neutrino mass matrix as:

$$\mathcal{M}_N = \begin{bmatrix} \mathrm{M}_{\chi^0} & m^T \\ m & 0 \end{bmatrix} \tag{12}$$

where M_{χ^0} is given by eq. (7) and m reads as:

$$m = \begin{bmatrix} -\frac{1}{2}\tilde{g}' c'_\beta a_1 v & \frac{1}{2}\tilde{g} c_\beta a_1 v & 0 & \epsilon_1 \\ -\frac{1}{2}\tilde{g}' c'_\beta a_2 v & \frac{1}{2}\tilde{g} c_\beta a_2 v & 0 & \epsilon_2 \\ -\frac{1}{2}\tilde{g}' c'_\beta a_3 v & \frac{1}{2}\tilde{g} c_\beta a_3 v & 0 & \epsilon_3 \end{bmatrix} \tag{13}$$

We define the parameters $\lambda_i \equiv a_i\mu + \epsilon_i$, which are related to the traditional BRpV parameters Λ_i [13] by $\Lambda_i = \lambda_i v_d$. Integrating out the neutralinos, we find that the neutrino mass matrix is given by:

$$\mathbf{M}_\nu^{eff} = -m\,\mathrm{M}_{\chi^0}^{-1}\,m^T = \frac{v^2}{4\det M_{\chi^0}}\left(M_1\tilde{g}^2 c_\beta^2 + M_2\tilde{g'}^2 c'^2_\beta\right)\begin{bmatrix} \lambda_1^2 & \lambda_1\lambda_2 & \lambda_1\lambda_3 \\ \lambda_2\lambda_1 & \lambda_2^2 & \lambda_2\lambda_3 \\ \lambda_3\lambda_1 & \lambda_3\lambda_2 & \lambda_3^2 \end{bmatrix} \tag{14}$$

where the determinant of the neutralino mass matrix is:

$$\det M_{\chi^0} = -\mu^2 M_1 M_2 + \frac{1}{2}v^2\mu\left(M_1\tilde{g}^2 s_\beta c_\beta + M_2\tilde{g'}^2 s'_\beta c'_\beta\right) + \tfrac{1}{16}\tilde{g}^2\tilde{g'}^2 v^4\left(s'_\beta c_\beta - s_\beta c'_\beta\right)^2 \tag{15}$$

Notice that the effective neutrino mass matrix $\mathbf{M}_\nu^{eff}$ has only one eigenvalue different from zero. Therefore, as in the case of R-parity violation in the MSSM with bilinear terms, at tree level only one neutrino is massive. Therefore, we have to investigate all possible one loop contributions to the neutrino mass matrix which help us to solve the atmospheric and solar neutrino problems. It has been argued in the literature that using the bilinear terms it is not possible to explain the neutrino masses and mixing. However, as we will show in the next sections, once we include the one-loop contributions to the neutrino mass matrix there is no problem to explain the neutrino properties.

A. One-loop Corrections to the Neutrino Mass Matrix

The one loop corrections are crucial for the correct characterization of neutrino phenomena. In the MSSM usually the most important one-loop contributions to the neutrino mass matrix are the bottom squarks, charginos, and neutralinos contributions. In Split Supersymmetry all scalars, except for one light Higgs boson, are super heavy. Therefore in this case the only potentially important contributions are charginos and neutralinos together with W, Z, and light Higgs inside the loop. We will show that Z

and W loops (see Appendix A) are just a small renormalization of the tree-level contribution, and that the crucial loops include the light Higgs and the neutralinos.

In general, the one loop contributions to the neutrino mass matrix can be written as [6]:

$$\Delta M_\nu^{ij} = \Pi_{ij}(0) = -\frac{1}{16\pi^2}\sum_{f,b} G_{ijfb} m_f B_0(0; m_f^2, m_b^2) \tag{16}$$

where the sum is over the fermions (f) and the bosons (b) inside the loop, m_f is the fermion mass, and G_{ijfb} is related to the couplings between the neutrinos and the fermions and bosons inside the loop. Once the smallness of the ϵ_i and λ_i parameters is taken into account, each contribution can be expressed in the form

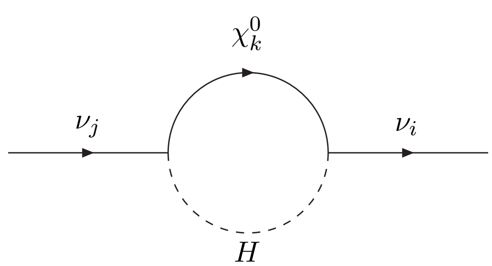
$$\Delta\Pi_{ij} = A^{(1)}\lambda_i\lambda_j + B^{(1)}(\epsilon_i\lambda_j + \epsilon_j\lambda_i) + C^{(1)}\epsilon_i\epsilon_j \tag{17}$$

with $A^{(1)}$, $B^{(1)}$, and $C^{(1)}$ parameters independent of ϵ_i and λ_i, but dependent on the other SUSY parameters. The super-index (1) refers to the one-loop contribution. The tree-level neutrino mass matrix in eq. (14) has the form $\mathbf{M}_{\nu\ ij}^{eff} = A^{(0)}\lambda_i\lambda_j$ with

$$A^{(0)} = \frac{v^2}{4\det M_{\chi^0}}\left(M_1\tilde{g}^2 c_\beta^2 + M_2\tilde{g'}^2 c'^2_\beta\right) \tag{18}$$

and we define the one-loop corrected parameters $A = A^{(0)} + A^{(1)}$, $B = B^{(1)}$, and $C = C^{(1)}$.

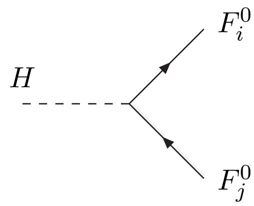
In the MSSM with BRpV the neutral Higgs bosons mix with the sneutrinos forming two sets of 5 scalars and 5 pseudo-scalars. Nevertheless, in Split supersymmetry, all the sneutrinos are extremely heavy and decouple from the light Higgs boson H. In addition, the heavy Higgs boson also has a very large mass, leaving the light Higgs as the only neutral scalar able to contribute to the neutrino masses. This contribution is represented by the following Feynman graph,



which is proportional to the neutralino mass $m_{\chi_k^0}$. Here χ_k^0 and H are the neutralino and Higgs mass eigenstates, but the graph is calculated in the basis where ν_i are not mass eigenstates. The fields ν_i are the neutrino fields associated to the effective mass matrix given in eq. (14). This contribution to eq. (16) proceeds with the coupling [6]

$$G_{ijk}^h = \frac{1}{2}(O_{Ljk}^{nnh}O_{Lki}^{nnh} + O_{Rjk}^{nnh}O_{Rki}^{nnh}) \tag{19}$$

where the relevant vertex is:



$$= i\left[O^{nnh}_{Lij}\frac{(1-\gamma_5)}{2} + O^{nnh}_{Rij}\frac{(1+\gamma_5)}{2}\right]$$

Here F_i^0 are the seven eigenvectors linear combination of the higgsinos, gauginos, and neutrinos. The O_L and O_R couplings are given by:

$$O^{nnh}_{Lij} = (O^{nnh}_{Rji})^* \,, \qquad O^{nnh}_{Rij} = S_{ij} - Q_{ij} \tag{20}$$

with

$$\begin{aligned} S_{ij} &= \tfrac{1}{2}\left[\mathcal{N}_{i4}\left(\tilde{g}s_\beta\mathcal{N}_{j2} - \tilde{g}'s'_\beta\mathcal{N}_{j1}\right) + \mathcal{N}_{j4}\left(\tilde{g}s_\beta\mathcal{N}_{i2} - \tilde{g}'s'_\beta\mathcal{N}_{i1}\right)\right] \\ Q_{ij} &= \tfrac{1}{2}\left[\mathcal{N}_{i3}\left(\tilde{g}c_\beta\mathcal{N}_{j2} - \tilde{g}'c'_\beta\mathcal{N}_{j1}\right) + \mathcal{N}_{j3}\left(\tilde{g}c_\beta\mathcal{N}_{i2} - \tilde{g}'c'_\beta\mathcal{N}_{i1}\right)\right] \end{aligned} \tag{21}$$

where the difference with the MSSM couplings given in [14] lies in the fact that in our case $\mathcal{N}$ is a 7×7 matrix, and the Higgs mixing angle has been replaced by $\alpha = \beta - \pi/2$, valid in the decoupling limit [15].

Using the approximation for the matrix $\mathcal{N}$ from Appendix A, we obtain:

$$\begin{aligned} O^{\nu\chi h}_{Rik} \approx & -\tfrac{1}{2}\left[\left(\tilde{g}s_\beta N_{k2} - \tilde{g}'s'_\beta N_{k1}\right)\xi_{i4} + N_{k4}\left(\tilde{g}s_\beta\,\xi_{i2} - \tilde{g}'s'_\beta\,\xi_{i1}\right)\right] \\ & +\tfrac{1}{2}\left[\left(\tilde{g}c_\beta N_{k2} - \tilde{g}'c'_\beta N_{k1}\right)\xi_{i3} + N_{k3}\left(\tilde{g}c_\beta\,\xi_{i2} - \tilde{g}'c'_\beta\,\xi_{i1}\right)\right] \end{aligned} \tag{22}$$

Notice that the presence of the term ξ_{i3} implies that the contribution of the light Higgs boson has the form:

$$\Delta\Pi^h_{ij} = A^h\lambda_i\lambda_j + B^h(\lambda_i\epsilon_j + \lambda_j\epsilon_i) + C^h\epsilon_i\epsilon_j \tag{23}$$

breaking the symmetry of the neutrino mass matrix at tree level. Explicitly, the graph above eq. (19) is:

$$\Delta\Pi^h_{ij} = -\frac{1}{64\pi^2}\sum_{k=1}^{4}(E_k\lambda_i + F_k\epsilon_i)(E_k\lambda_j + F_k\epsilon_j)\, m_{\chi^0_k} B_0(0; m^2_{\chi^0_k}, m^2_H) \tag{24}$$

with

$$\begin{aligned} E_k &= -\left(\tilde{g}s_\beta N_{k2} - \tilde{g}'s'_\beta N_{k1}\right)\xi_4 - N_{k4}\left(\tilde{g}s_\beta\,\xi_2 - \tilde{g}'s'_\beta\,\xi_1\right) \\ & \quad + \left(\tilde{g}c_\beta N_{k2} - \tilde{g}'c'_\beta N_{k1}\right)\xi_3 + N_{k3}\left(\tilde{g}c_\beta\,\xi_2 - \tilde{g}'c'_\beta\,\xi_1\right) \\ F_k &= -\frac{1}{\mu}\left(\tilde{g}c_\beta N_{k2} - \tilde{g}'c'_\beta N_{k1}\right) \end{aligned} \tag{25}$$

where we have used the notation in Appendix A. We work in the Feynman gauge, and the loop with a Goldstone boson G must be included, as described in Appendix A. This contribution $\Delta\Pi^G_{ij}$ can be obtained from $\Delta\Pi^h_{ij}$ in eq. (24) doing the following: i) change the overall sign, ii) change m^2_H by m^2_Z inside the Veltman function, and iii) change the sign of s_β and s'_β in the definition of E_k. When the Goldstone boson contribution is added to the Higgs contribution, a partial cancellation occurs, making the value of the C parameter smaller than in Sugra, where other contributions are present (see for example [16]). In appendix B we explain how this cancellation is more severe in the case of two light Higgs doublets, also showing how the CP-even and the CP-odd contributions cancel each other in this situation, found in Sugra for example.

It is useful to understand the origin of the C^h coefficient in eq. (23), which is one of the two terms that break the symmetry of the tree-level neutrino mass matrix in eq. (14). This coefficient is proportional to F^2_k, given in eq. (25), from which three kind of terms appear: one proportional to $(\tilde{g}c_\beta N_{k2})^2$, one proportional to $(\tilde{g}'c'_\beta N_{k1})^2$ and a mixing. In the first term $\tilde{g}c_\beta = \tilde{g}_d$ according to our notation, and corresponds to the vertex $H\widetilde{W}\widetilde{H}_d$ given in eq. (1). This means that this contribution comes from a Higgs boson H plus a wino $\widetilde{W}$ inside the loop with a higgsino $\widetilde{H}_d$ in the external legs. Therefore, the two vertices account for the factor $\tilde{g}^2c^2_\beta$, and the projection of the neutralino χ^0_k (the mass eigenstate inside the loop) into the wino accounts for the factor N^2_{k2}. Finally, in the external legs the higgsino $\widetilde{H}_d$ mixes with the neutrinos introducing the factor $\epsilon_i\epsilon_j/\mu^2$. The term proportional to N^2_{k1} is analogous to the above one replacing the winos by binos in the two vertices. The mixing term can be understood in a similar way, as well as the B^h coefficient. Therefore, the loops with a Higgs boson and gauginos inside it form the most important contribution to the neutrino mass matrix, which relaxes the masslessness of the solar neutrino scale.

B. Numerical Results

For our model to be viable it must generate neutrino masses and mixing angles compatible with experimental data. Global fits of the data allow to extract from it the two neutrino mass squared differences Δm^2_{32} and Δm^2_{12}, and the three mixing angles θ_{23}, θ_{12} and θ_{13}. For the atmospheric and solar parameters we use [17],

$$\begin{aligned} &1.4\times10^{-3} < \Delta m^2_{32} < 3.3\times10^{-3}\,\mathrm{eV}^2\,, \qquad && 0.52 < \tan^2\theta_{23} < 2.1\,, \\ &7.2\times10^{-5} < \Delta m^2_{21} < 9.1\times10^{-5}\,\mathrm{eV}^2\,, \qquad && 0.30 < \tan^2\theta_{12} < 0.61\,, \end{aligned} \tag{26}$$

which are completed with the upper bounds for the neutrinoless double beta decay mass parameter and the reactor angle,

$$m_{ee} < 0.84\,\mathrm{eV}\,, \qquad \tan^2\theta_{13} < 0.049\,, \tag{27}$$

with all bounds valid at 3σ.

As a supersymmetric working scenario we choose the values indicated in Table I. We need to specify the input parameters for the neutralino and chargino mass matrices: the gaugino masses $M_1 = 50$ GeV and $M_2 = 300$ GeV, the higgsino mass parameter $\mu = 200$ GeV, and $\tan\beta = 50$ (for the experimental bounds on the neutralino masses see [18]). In addition, the Higgs mass is taken to be $m_H = 120$ GeV. Part of the neutralino and chargino spectrum is also indicated in Table I. The LSP is the lightest neutralino with a mass $m_{\chi^0_1} = 47$ GeV, which is unstable.

Within this supersymmetric scenario we found a proof of existence solution which satisfy all neutrino data in Eqs. (26) and (27). The predictions of this scenario are given in Table II. This solution is characterized by a non-maximal pair of atmospheric and solar mixing angles, with $\tan^2\theta_{sol} < 1 < \tan^2\theta_{atm}$, and by atmospheric and solar mass squared differences relatively centered in their corresponding allowed ranges.

The one-loop corrected A, B, and C parameters that define the neutrino mass matrix, as defined below eq. (18), are $A = -620\,,\ B = -1.4\,$, and $C = 0.25$, all of them in units of eV/GeV^2. The A parameter is almost entirely generated at tree-level, $A^{(0)} = -604$ eV/GeV^2, with the contributions from W, Z, G, and H loops to A being minimal. The importance of the one-loop generated parameters is enhanced because the ϵ_i are larger than the λ_i parameters. Neglecting ϵ_2 and λ_1 we have,

$$\mathbf{M}_\nu^{eff} \approx \begin{bmatrix} 0 & 0 & 0 \\ 0 & A\lambda_2^2 & A\lambda_2\lambda_3 \\ 0 & A\lambda_2\lambda_3 & A\lambda_3^2 \end{bmatrix} + \begin{bmatrix} 0 & B\epsilon_1\lambda_2 & B\epsilon_1\lambda_3 \\ B\epsilon_1\lambda_2 & 0 & B\epsilon_3\lambda_2 \\ B\epsilon_1\lambda_3 & B\epsilon_3\lambda_2 & 2B\epsilon_3\lambda_3 \end{bmatrix} + \begin{bmatrix} C\epsilon_1^2 & 0 & C\epsilon_1\epsilon_3 \\ 0 & 0 & 0 \\ C\epsilon_1\epsilon_3 & 0 & C\epsilon_3^2 \end{bmatrix} \tag{28}$$

This mass matrix has an exact zero eigenvalue, and this agrees with the complete numerical calculation. If we treat the terms proportional to B and C as perturbations, the two massive neutrinos have mass given by,

$$m_{\nu_3} = A|\vec{\lambda}|^2 + 2B(\vec{\epsilon}\cdot\vec{\lambda}) + C\frac{(\vec{\epsilon}\cdot\vec{\lambda})^2}{|\vec{\lambda}|^2} \tag{29}$$

$$m_{\nu_2} = C\frac{|\vec{\lambda}\times(\vec{\epsilon}\times\vec{\lambda})|^2}{|\vec{\lambda}|^4} \tag{30}$$

where $\vec{\epsilon} = (\epsilon_1, \epsilon_2, \epsilon_3)$ and $\vec{\lambda} = (\lambda_1, \lambda_2, \lambda_3)$. We see that the solar mass squared is generated in first approximation only by the C term. The solar mass squared difference is $\Delta m^2_{sol} = m^2_{\nu_2}$, and the approximation in eq. (30) gives a value of $6\times 10^{-5}\,\mathrm{eV}^2$, very close to the exact value in Table II. Similarly, the atmospherics mass squared difference, given by $\Delta m^2_{atm} = m^3_{\nu_2} - m^3_{\nu_2}$, using the approximation in eq. (30) gives $2\times 10^{-3}\,\mathrm{eV}^2$ also in agreement with the exact numerical answer in Table II.

In Fig. 1 we take the Split-BRpV scenario defined by the parameters in Tables I and II and vary the wino mass M_2 and the higgsino mass μ, looking for the points in the $M_2 - \mu$ plane that satisfy the neutrino parameter bounds given in Eqs. (26) and (27). The result is shown as the shaded (yellow) region in Fig. 1. The boundaries of this region are the upper and lower limits on the atmospheric and solar mass squared differences. For a fix value of the wino mass, both mass squared differences decrease once we increase the higgsino mass. Notice that μ acts mainly as a scale factor for the ϵ_i parameters, as seen in eq. (25). The tangent of the atmospheric angle increases with growing M_2, and depends negligibly on μ. On the other hand, the dependence of the solar angle on both of these parameters is negligible.

The solution in Fig. 1 shows that the atmospheric mass squared difference decreases with M_2 at fixed higgsino mass. The dependence of A on the gaugino mass M_2 is such that A is negative and its absolute value decreases with increasing M_2. Since A is mainly given by tree-level contributions, this behavior can be understood from eq. (18). The sign of A comes from the first term in eq. (15) which is the dominant one.

In this way we show that once the one-loop corrections (particularly the Higgs contribution) to the neutrino mass matrix are included in the analysis it is possible to explain all neutrino properties. Our main conclusion is that even if the SUSY scale (the sfermion masses) is very large we can use the bilinear interactions in order to describe all neutrino properties. We hope that those results will motivate new studies in the context of supersymmetric scenarios with large sfermion masses and R-parity violation.

III. SUMMARY

We have studied in detail the possibility to describe the neutrino masses and mixing angles in the context of split supersymmetric scenarios. We have considered all relevant contributions to the neutrino mass matrix up to one-loop level coming from the R-parity violating interactions, showing the importance of the Higgs one-loop correction. Contrary to previous studies, we conclude that it is possible to explain the neutrino properties using the bilinear R-Parity violating interactions.

Acknowledgments

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TABLE I: Supersymmetric parameters and part of the spectrum in Split Susy benchmark.

Parameter	Benchmark A	Units
M_1	50	GeV
M_2	300	GeV
μ	200	GeV
m_h	120	GeV
$\tan\beta$	50	
$m_{\chi_1^0}$	47	GeV
$m_{\chi_2^0}$	181	GeV
$m_{\chi_1^+}$	180	GeV
$m_{\chi_2^+}$	333	GeV

TABLE II: BRpV and neutrino parameters in Split Susy benchmark.

Parameter	Benchmark	Units
ϵ_1	-0.107	GeV
ϵ_2	-0.001	GeV
ϵ_3	0.217	GeV
λ_1	2×10^{-5}	GeV
λ_2	0.0063	GeV
λ_3	-0.0073	GeV
Δm^2_{atm}	2.34×10^{-3}	eV^2
Δm^2_{sol}	8.37×10^{-5}	eV^2
$\tan^2\theta_{atm}$	1.22	-
$\tan^2\theta_{sol}$	0.56	-
$\tan^2\theta_{13}$	0.0083	-
m_{ee}	0.0037	eV

APPENDIX A: GAUGE BOSON LOOPS

In this appendix we show the properties of the gauge boson one-loop contributions to the neutrino mass matrix.

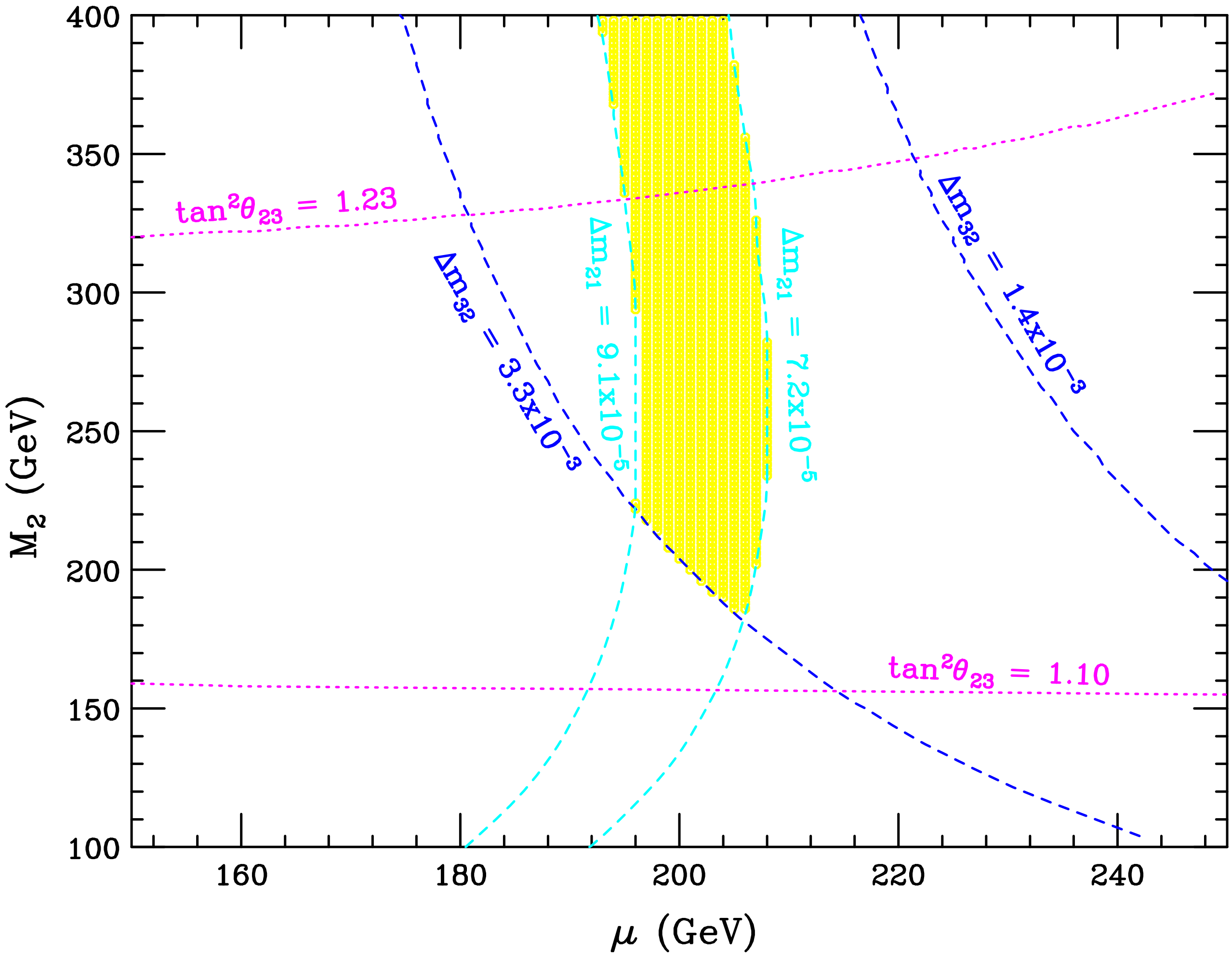


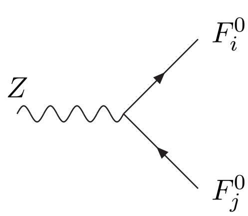
FIG. 1: *Solutions for neutrino physics with Split Supersymmetry benchmark values for ϵ and Λ in the $M_2 - \mu$ plane .*

1. Z and neutral Goldstone boson loops

In Z loops the fermionic sum in eq. (16) is over neutral fermions F_k^0, of which only the neutralinos are relevant. There is no bosonic sum since only Z contributes. The coupling G_{ijk}^Z is equal to

$$G_{ijk}^Z = -2(O_{Ljk}^{nnz}O_{Rki}^{nnz} + O_{Rjk}^{nnz}O_{Lki}^{nnz}) \tag{A1}$$

where the coupling of a Z boson to two neutral fermions is



$$= i\left[O_{Lij}^{nnz}\frac{(1-\gamma_5)}{2} + O_{Rij}^{nnz}\frac{(1+\gamma_5)}{2}\right]$$

with

$$O^{nnz}_{Lij} = -(O^{nnz}_{Rij})^* \,, \qquad O^{nnz}_{Rij} = -\frac{g}{2c_W}\left(\mathcal{N}^*_{i4}\mathcal{N}_{j4} - \mathcal{N}^*_{i3}\mathcal{N}_{j3} - \sum_{a=1}^{3}\mathcal{N}^*_{ia+4}\mathcal{N}_{ja+4}\right) \tag{A2}$$

The matrix $\mathcal{N}$ diagonalizes the 7×7 neutrino/neutralino mass matrix, giving non-negative eigenvalues. Without including the final rotation on the neutrino sector, it can be approximated in the following way [6]:

$$\mathcal{N} \approx \begin{bmatrix} N & N\xi^T \\ -\xi & 1 \end{bmatrix} \tag{A3}$$

where N diagonalizes the 4×4 neutralino mass sub-matrix. The parameters ξ are defined by

$$\begin{aligned}
\xi_{i1} &= \frac{\tilde{g}'v}{2\det M_{\chi^0}}\left[c'_\beta \mu M_2 + \frac{1}{4}\tilde{g}^2 c_\beta v^2\left(s'_\beta c_\beta - s_\beta c'_\beta\right)\right]\lambda_i \\
\xi_{i2} &= -\frac{\tilde{g}v}{2\det M_{\chi^0}}\left[c_\beta \mu M_1 - \frac{1}{4}\tilde{g'}^2 c'_\beta v^2\left(s'_\beta c_\beta - s_\beta c'_\beta\right)\right]\lambda_i \\
\xi_{i3} &= \frac{v^2}{4\det M_{\chi^0}}\left[M_1\tilde{g}^2 s_\beta c_\beta + M_2\tilde{g'}^2 s'_\beta c'_\beta + \frac{1}{4}\tilde{g}^2\tilde{g'}^2\frac{v^2}{\mu}\left(s'_\beta c_\beta - s_\beta c'_\beta\right)^2\right]\lambda_i - \frac{\epsilon_i}{\mu} \\
\xi_{i4} &= -\frac{v^2}{4\det M_{\chi^0}}\left(M_1\tilde{g}^2 c^2_\beta + M_2\tilde{g'}^2 {c'}^2_\beta\right)\lambda_i
\end{aligned} \tag{A4}$$

For notational brevity we define the ξ_i parameters as: $\lambda_i\xi_1 = \xi_{i1}$, $\lambda_i\xi_2 = \xi_{i2}$, $\lambda_i\xi_3 - \epsilon_i/\mu = \xi_{i3}$, and $\lambda_i\xi_4 = \xi_{i4}$. The couplings in eq. (A2) can be approximated with the help of eq. (A3) to

$$O^{\nu\chi z}_{Rik} \approx \frac{g}{2c_W}\left(2N_{k4}\xi_{i4} + N_{k1}\xi_{i1} + N_{k2}\xi_{i2}\right) \tag{A5}$$

where i labels the three neutrinos and k labels the four neutralinos. Considering eq. (A4) we conclude,

$$\Delta\Pi^Z_{ij} = A^Z\lambda_i\lambda_j \tag{A6}$$

with

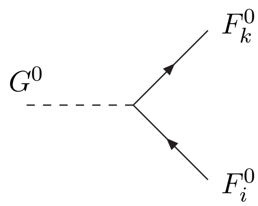
$$A^Z = -\frac{g^2}{16\pi^2 c^2_W}\sum_{k=1}^{4}\left(2N_{k4}\xi_4 + N_{k1}\xi_1 + N_{k2}\xi_2\right)^2 m_{\chi^0_k} B_0(0; m^2_{\chi^0_k}, m^2_Z) \tag{A7}$$

This contribution is only a renormalization of the tree level mass matrix which it does not break its symmetry, i.e., it does not generate mass to all neutrinos.

For loops containing pseudoscalars, the sum in eq. (16) is over neutral fermions F^0_k and neutral pseudoscalars P^0_k. Nevertheless, only neutralinos and the neutral Goldstone boson are relevant. The coupling G^0_{ijk} for the neutral Goldstone boson is,

$$G^0_{ijk} = -\frac{1}{2}(O^{nng}_{Ljk}O^{nng}_{Lik} + O^{nng}_{Rjk}O^{nng}_{Rik}) \tag{A8}$$

The corresponding vertex is



$$= i\left[O^{nng}_{Lik}\,\tfrac{(1-\gamma_5)}{2} + O^{nng}_{Rik}\,\tfrac{(1+\gamma_5)}{2}\right]$$

with

$$O^{nng}_{Lik} = -(O^{nng}_{Rki})^*\,, \qquad O^{nng}_{Rik} = -S_{ik} - Q_{ik} \tag{A9}$$

where S and Q are defined in eq. (21). The neutral Goldstone boson contribution (in the Feynman gauge) to the neutrino mass matrix is

$$\Delta\Pi^{G^0}_{ij} \;=\; \frac{1}{64\pi^2}\sum_{k=1}^{4}\left(E^0_k\lambda_i + F^0_k\epsilon_i\right)\left(E^0_k\lambda_j + F^0_k\epsilon_j\right) m_{\chi^0_k} B_0(0; m^2_{\chi^0_k}, m^2_Z) \tag{A10}$$

with

$$\begin{aligned} E^0_k \;&=\; \left(\tilde g s_\beta N_{k2} - \tilde g' s'_\beta N_{k1}\right)\xi_4 + N_{k4}\left(\tilde g s_\beta\,\xi_2 - \tilde g' s'_\beta\,\xi_1\right) \\ &\quad + \left(\tilde g c_\beta N_{k2} - \tilde g' c'_\beta N_{k1}\right)\xi_3 + N_{k3}\left(\tilde g c_\beta\,\xi_2 - \tilde g' c'_\beta\,\xi_1\right) \\ F^0_k \;&=\; -\frac{1}{\mu}\left(\tilde g c_\beta N_{k2} - \tilde g' c'_\beta N_{k1}\right) \end{aligned} \tag{A11}$$

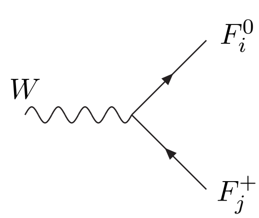
As explain before, this contribution tends to cancel the $\epsilon_i\epsilon_j$ term from the Higgs contribution, but this cancellation is only partial and do not spoil the solution for neutrino masses and oscillations.

2. W and charged Goldstone boson loops

In W loops the fermionic sum in eq. (16) is over charged fermions F^+_k, of which only the charginos are relevant. There is no bosonic sum since only W contributes. The coupling G^W_{ijk} is equal to

$$G^W_{ijk} = -4(O^{ncw}_{Ljk}O^{ncw}_{Rik} + O^{ncw}_{Rjk}O^{ncw}_{Lik}) \tag{A12}$$

where the coupling of a W boson to two fermions is



$$= i\left[O^{ncw}_{Lij}\,\tfrac{(1-\gamma_5)}{2} + O^{ncw}_{Rij}\,\tfrac{(1+\gamma_5)}{2}\right]$$

with

$$O^{ncw}_{Lij} = -g\left(\mathcal{N}^*_{i2}\mathcal{U}_{j1} + \frac{1}{\sqrt{2}}\mathcal{N}^*_{i3}\mathcal{U}_{j2} + \frac{1}{\sqrt{2}}\sum_{a=1}^{3}\mathcal{N}^*_{ia+4}\mathcal{U}_{ja+2}\right)$$
$$O^{ncw}_{Rij} = -g\left(\mathcal{N}_{i2}\mathcal{V}^*_{j1} - \frac{1}{\sqrt{2}}\mathcal{N}_{i4}\mathcal{V}^*_{j2}\right) \tag{A13}$$

The $\mathcal{U}$ and $\mathcal{V}$ matrices diagonalize the 5×5 chargino/charged lepton mass matrix, and can be approximated to [6]

$$\mathcal{U} \approx \begin{bmatrix} U & U\xi_L^T \\ -\xi_L & 1 \end{bmatrix}, \qquad \mathcal{V} \approx \begin{bmatrix} V & 0 \\ 0 & 1 \end{bmatrix} \tag{A14}$$

where U and V diagonalize the 2×2 chargino sub-matrix. The parameters ξ_L are

$$\xi_L^{i1} = \frac{\tilde{g}vc_\beta}{\sqrt{2}\det M_{\chi^+}}\lambda_i\,, \qquad \xi_L^{i2} = -\frac{\tilde{g}^2v^2s_\beta c_\beta}{2\mu\det M_{\chi^+}}\lambda_i - \frac{\epsilon_i}{\mu}\,, \tag{A15}$$

with

$$\det M_{\chi^+} = \mu M_2 - \frac{1}{2}\tilde{g}^2v^2s_\beta c_\beta \tag{A16}$$

and similarly to what we did in the previous subsection, we define the parameters ξ_j^L, $j=1,2$, with the relations: $\xi_L^{i1}=\xi_1^L\lambda_i$ and $\xi_L^{i2}=\xi_2^L\lambda_i-\epsilon_i/\mu$. The couplings in eq. (A13) can be approximated to

$$O^{\nu\chi w}_{Rij} \approx g\left(V^*_{j1}\xi_{i2} - \frac{1}{\sqrt{2}}V^*_{j2}\xi_{i4}\right)$$
$$O^{\nu\chi w}_{Lij} \approx g\left(U_{j1}\xi_{i2} - \frac{1}{\sqrt{2}}U_{j2}\left[\xi_L^{i2}-\xi_{i3}\right] + \frac{1}{\sqrt{2}}U_{j1}\xi_L^{i1}\right) \tag{A17}$$

where i labels the three neutrinos and j labels the two charginos. Similarly to what happened with the Z contributions, the W contribution depends only on the λ_i:

$$\Delta\Pi^W_{ij} = A^W\lambda_i\lambda_j \tag{A18}$$

with

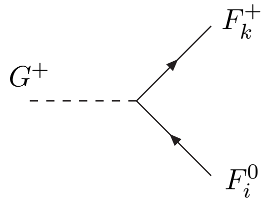
$$A^W = \frac{g^2}{2\pi^2}\sum_{k=1}^{2}\left[U_{k1}\xi_2 - \frac{U_{k2}}{\sqrt{2}}\left(\xi_2^L-\xi_3\right) + \frac{U_{k1}}{\sqrt{2}}\xi_1^L\right]\left(V_{k1}\xi_2 - \frac{V_{k2}}{\sqrt{2}}\xi_4\right)m_{\chi_k^+}B_0(0;m^2_{\chi_k^+},m_W^2) \tag{A19}$$

Adding to the tree level contribution without changing the symmetry. Therefore the W and Z loops do not help us to generate mass to all neutrinos.

For loops containing charged scalars, the sum in eq. (16) is over charged fermions F_k^+ and charged scalars S_k^+, and among them only the charginos and the charged Goldstone boson are relevant. The coupling G^+_{ijk} for the charged Goldstone boson is,

$$G^+_{ijk} = (O^{ncg}_{Ljk}O^{ncg}_{Rik} + O^{ncg}_{Rjk}O^{ncg}_{Lik}) \tag{A20}$$

where the relevant vertex is:



$$= i\left[O^{ncg}_{Lik}\tfrac{(1-\gamma_5)}{2} + O^{ncg}_{Rik}\tfrac{(1+\gamma_5)}{2}\right]$$

with

$$\begin{aligned} O^{ncg}_{Lik} &= -\tilde{g}s_\beta\mathcal{N}_{i4}\mathcal{V}_{k1} - \frac{1}{\sqrt{2}}(\tilde{g}s_\beta\mathcal{N}_{i2} + \tilde{g}'s'_\beta\mathcal{N}_{i1})\mathcal{V}_{k2} \\ O^{ncg}_{Rik} &= \tilde{g}c_\beta\mathcal{N}_{i3}\mathcal{U}_{k1} - \frac{1}{\sqrt{2}}(\tilde{g}c_\beta\mathcal{N}_{i2} + \tilde{g}'c'_\beta\mathcal{N}_{i1})\mathcal{U}_{k2} \end{aligned} \tag{A21}$$

Using the approximations in eq. (A3) and (A14) we find,

$$\begin{aligned} O^{\nu\chi g}_{Lik} &= \tilde{g}s_\beta\xi_{i4}V_{k1} + \frac{1}{\sqrt{2}}(\tilde{g}s_\beta\xi_{i2} + \tilde{g}'s'_\beta\xi_{i1})V_{k2} \\ O^{\nu\chi g}_{Rik} &= -\tilde{g}c_\beta\xi_{i3}U_{k1} + \frac{1}{\sqrt{2}}(\tilde{g}c_\beta\xi_{i2} + \tilde{g}'c'_\beta\xi_{i1})U_{k2} \end{aligned} \tag{A22}$$

In this way, the charged Goldstone boson contribution to the neutrino mass matrix can be written as,

$$\Delta\Pi^{G^+}_{ij} = -\frac{1}{32\pi^2}\sum_{k=1}^{2}\left[2E^+_{Lk}E^+_{Rk}\lambda_i\lambda_j + E^+_{Lk}F^+_{Rk}\left(\epsilon_i\lambda_j + \epsilon_j\lambda_i\right)\right] m_{\chi^+_k} B_0(0; m^2_{\chi^+_k}, m^2_W) \tag{A23}$$

where

$$\begin{aligned} E^+_{Lk} &= \tilde{g}s_\beta\xi_4V_{k1} + \frac{1}{\sqrt{2}}(\tilde{g}s_\beta\xi_2 + \tilde{g}'s'_\beta\xi_1)V_{k2} \\ E^+_{Rk} &= -\tilde{g}c_\beta\xi_3U_{k1} + \frac{1}{\sqrt{2}}(\tilde{g}c_\beta\xi_2 + \tilde{g}'c'_\beta\xi_1)U_{k2} \\ F^+_{Rk} &= \tilde{g}c_\beta U_{k1}/\mu \end{aligned} \tag{A24}$$

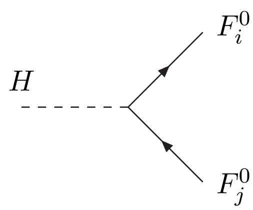
Notice that the charged Goldstone boson does not generates a $\epsilon_i\epsilon_j$ term.

APPENDIX B: EFFECT OF THE DECOUPLING

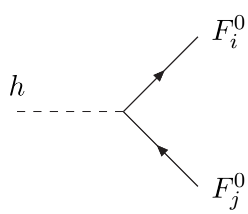
In this appendix we study in detail the effect of the decoupling of the scalar particles, and compare with previous results in the literature.

1. Decoupling Limit

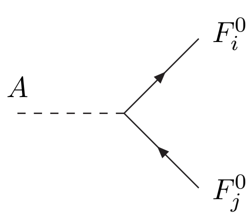
Here we compare our results with the case when the two Higgs doublets are light. We assume all sneutrinos are very heavy and their mixing with the Higgs bosons are negligible in the calculation that follows. Our main contribution comes from Higgs bosons and gauginos inside the loop, and a higgsino mixing with neutrinos in the external leg. The relevant vertex is the one with a Higgs boson and two neutralinos, given in ref. [6] for the case of MSSM-BRpV, and in [14] and [19] for the case of the MSSM. If for simplicity we assume all parameters in the neutralino mass matrix are real and that its eigenvalues are positive, then the relevant Feynman rules are reduced to (in this appendix we call h the light Higgs boson and H the heavy Higgs boson),



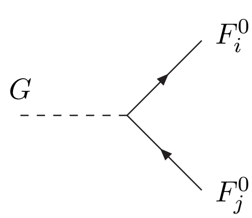
$$= i\, O_{ij}^{nnH} = -i\,(Q_{ij}^{nn} c_\alpha - S_{ij}^{nn} s_\alpha)$$



$$= i\, O_{ij}^{nnh} = i\,(Q_{ij}^{nn} s_\alpha + S_{ij}^{nn} c_\alpha)$$



$$= O_{ij}^{nnA}\gamma_5 = (Q_{ij}^{nn} s_\beta - S_{ij}^{nn} c_\beta)\gamma_5$$



$$= O_{ij}^{nnG}\gamma_5 = -(Q_{ij}^{nn} c_\beta + S_{ij}^{nn} s_\beta)\gamma_5$$

with

$$\begin{aligned} Q_{ij}^{nn} &= \tfrac{1}{2}\left[\mathcal{N}_{i3}(g\mathcal{N}_{j2} - g'\mathcal{N}_{j1}) + \mathcal{N}_{j3}(g\mathcal{N}_{i2} - g'\mathcal{N}_{i1})\right] \\ S_{ij}^{nn} &= \tfrac{1}{2}\left[\mathcal{N}_{i4}(g\mathcal{N}_{j2} - g'\mathcal{N}_{j1}) + \mathcal{N}_{j4}(g\mathcal{N}_{i2} - g'\mathcal{N}_{i1})\right] \end{aligned} \tag{B1}$$

When one of the neutral fermions is a neutrino and the other is a neutralino, the above couplings can be approximated to

$$\begin{aligned} Q_{ik}^{\nu\chi} &= \tfrac{1}{2}\left[-\xi_{i3}(gN_{k2} - g'N_{k1}) - N_{k3}(g\xi_{i2} - g'\xi_{i1})\right] \\ S_{ik}^{\nu\chi} &= \tfrac{1}{2}\left[-\xi_{i4}(gN_{k2} - g'N_{k1}) - N_{k4}(g\xi_{i2} - g'\xi_{i1})\right] \end{aligned} \tag{B2}$$

The loops involving the k^{th} neutralino contribute with

$$\Delta\Pi_{ij}^{0} = -\frac{m_{\chi_k^0}}{16\pi^2}\left[O_{ik}^{\nu\chi H}O_{jk}^{\nu\chi H}B_0^{0kH} + O_{ik}^{\nu\chi h}O_{jk}^{\nu\chi h}B_0^{0kh} - O_{ik}^{\nu\chi A}O_{jk}^{\nu\chi A}B_0^{0kA} - O_{ik}^{\nu\chi G}O_{jk}^{\nu\chi G}B_0^{0kG}\right] \tag{B3}$$

In the limit where all scalar masses are the same, we can factor out the Veltman function and we find,

$$\begin{aligned} O_{ik}^{\nu\chi H}O_{jk}^{\nu\chi H} + O_{ik}^{\nu\chi h}O_{jk}^{\nu\chi h} &= (Q_{ik}^{\nu\chi}c_\alpha - S_{ik}^{\nu\chi}s_\alpha)(Q_{jk}^{\nu\chi}c_\alpha - S_{jk}^{\nu\chi}s_\alpha) + (Q_{ik}^{\nu\chi}s_\alpha + S_{ik}^{\nu\chi}c_\alpha)(Q_{jk}^{\nu\chi}s_\alpha + S_{jk}^{\nu\chi}c_\alpha) \\ &= Q_{ik}^{\nu\chi}Q_{jk}^{\nu\chi} + S_{ik}^{\nu\chi}S_{jk}^{\nu\chi} \end{aligned} \tag{B4}$$

and

$$\begin{aligned} O_{ik}^{\nu\chi A}O_{jk}^{\nu\chi A} + O_{ik}^{\nu\chi G}O_{jk}^{\nu\chi G} &= (Q_{ik}^{\nu\chi}s_\beta - S_{ik}^{\nu\chi}c_\beta)(Q_{jk}^{\nu\chi}s_\beta - S_{jk}^{\nu\chi}c_\beta) + (Q_{ik}^{\nu\chi}c_\beta + S_{ik}^{\nu\chi}s_\beta)(Q_{jk}^{\nu\chi}c_\beta + S_{jk}^{\nu\chi}s_\beta) \\ &= Q_{ik}^{\nu\chi}Q_{jk}^{\nu\chi} + S_{ik}^{\nu\chi}S_{jk}^{\nu\chi} \end{aligned} \tag{B5}$$

explaining the cancellation between the scalar and pseudoscalar loops in the MSSM with BRpV.

In the case of Split Supersymmetry H and A are decoupled and we replace $\sin\alpha = -\cos\beta$ and $\cos\alpha = \sin\beta$. The contribution from h and G is

$$\Delta\Pi_{ij}^{0} = -\frac{m_{\chi_k^0}}{16\pi^2}\left[O_{ik}^{\nu\chi h}O_{jk}^{\nu\chi h}B_0^{0kh} - O_{ik}^{\nu\chi G}O_{jk}^{\nu\chi G}B_0^{0kG}\right] \tag{B6}$$

In the limit where both states have the same mass the Veltman function can be factor out and the couplings satisfy,

$$\begin{aligned} O_{ik}^{\nu\chi h}O_{jk}^{\nu\chi h} - O_{ik}^{\nu\chi G}O_{jk}^{\nu\chi G} &= (Q_{ik}^{\nu\chi}c_\beta - S_{ik}^{\nu\chi}s_\beta)(Q_{jk}^{\nu\chi}c_\beta - S_{jk}^{\nu\chi}s_\beta) - (Q_{ik}^{\nu\chi}c_\beta + S_{ik}^{\nu\chi}s_\beta)(Q_{jk}^{\nu\chi}c_\beta + S_{jk}^{\nu\chi}s_\beta) \\ &= -2s_\beta c_\beta(Q_{ik}^{\nu\chi}S_{jk}^{\nu\chi} + Q_{jk}^{\nu\chi}S_{ik}^{\nu\chi}) \end{aligned} \tag{B7}$$

and the cancellation does not occur. The presence of the coupling $Q_{ik}^{\nu\chi}$ guarantees the presence of the term ϵ_i in the neutrino mass matrix, as can be seen from eq. (B2) and (A4). In this way, a B term is generated for the neutrino mass matrix as indicated in eq. (17), with a suppressed C term, breaking the tree level neutrino mass matrix symmetry, and thus generating a solar mass.

2. Sneutrino Higgs Mixing

In order to compare with previous results in the literature, we review here the mixing between Higgs bosons and sneutrino fields. The CP-even Higgs and sneutrino fields mix to form a set of five neutral mass eigenstates S_i^0, whose mass matrix is composed of the following blocks [6],

$$\mathbf{M}_{S^0}^2 = \begin{bmatrix} \mathbf{M}_{hh}^2 & \mathbf{M}_{h\widetilde{\nu}}^2 \\ \mathbf{M}_{h\widetilde{\nu}}^{2T} & \mathbf{M}_{\widetilde{\nu}\widetilde{\nu}}^2 \end{bmatrix} \tag{B8}$$

in the basis $(h_d, h_u, \tilde{\nu}_1, \tilde{\nu}_2, \tilde{\nu}_3)$. The higgs 2×2 sub-matrix is equal to,

$$\mathbf{M}_{hh}^2 = \begin{bmatrix} B_0\mu\frac{v_u}{v_d} + \frac{1}{4}g_Z^2 v_d^2 + \mu\vec{\epsilon}\cdot\frac{\vec{v}}{v_d} + \frac{t_d}{v_d} & -B_0\mu - \frac{1}{4}g_Z^2 v_d v_u \\ -B_0\mu - \frac{1}{4}g_Z^2 v_d v_u & B_0\mu\frac{v_d}{v_u} + \frac{1}{4}g_Z^2 v_u^2 - \vec{B}_\epsilon\cdot\frac{\vec{v}}{v_u} + \frac{t_u}{v_u} \end{bmatrix} \tag{B9}$$

where we call $g_Z^2 = g^2 + g'^2$, and in supergravity models we have $B_\epsilon^i = B_i\epsilon_i$. In this matrix we have eliminated the Higgs soft masses using the minimization conditions of the scalar potential (or tadpole equations) [6]. These Higgs tadpole equations at tree level are,

$$\begin{aligned} t_d &= \left(m_{H_d}^2 + \mu^2\right)v_d + v_d D - \mu\left(B_0 v_u + \vec{v}\cdot\vec{\epsilon}\right) \\ t_u &= -B_0\mu v_d + \left(m_{H_u}^2 + \mu^2\right)v_u - v_u D + \vec{v}\cdot\vec{B}_\epsilon + v_u\vec{\epsilon}^2 \end{aligned}$$

with $D = \frac{1}{8}(g^2 + g'^2)(\vec{v}^{\,2} + v_d^2 - v_u^2)$. At tree level, it is safe to set $t_u = t_d = 0$, and if we take the R-Parity conserving limit $\epsilon_i, v_i \to 0$, we can recognize the CP-even Higgs mass matrix of the MSSM. The 2×3 mixing sub-matrix is given by,

$$\mathbf{M}_{h\widetilde{\nu}}^2 = \begin{bmatrix} -\mu\epsilon_i + \frac{1}{4}g_Z^2 v_d v_i \\ B_\epsilon^i - \frac{1}{4}g_Z^2 v_u v_i \end{bmatrix} \tag{B10}$$

which vanishes in the R-Parity conserving limit. Finally, the sneutrino sub-matrix is given by,

$$\left(\mathbf{M}_{\widetilde{\nu}\widetilde{\nu}}^2\right)_{ij} = \left(M_{Li}^2 + D\right)\delta_{ij} + \frac{1}{4}g_Z^2 v_i v_j + \epsilon_i\epsilon_j \tag{B11}$$

where we have not yet used the corresponding tadpole equations, and we have assumed that the sneutrino soft mass matrix is diagonal. The sneutrino tadpole equations are given by,

$$t_i = v_i D + \epsilon_i\left(-\mu v_d + \vec{v}\cdot\vec{\epsilon}\right) + v_u B_\epsilon^i + v_i M_{Li}^2 \tag{B12}$$

It is clear from this equation that if the sneutrino vev's are zero, $\mu\epsilon_i = B_\epsilon^i v_u/v_d$, and therefore, the mixing between the up and down Higgs fields with the sneutrino fields are related by $M_{h_d\tilde{\nu}}^2 = -\tan\beta M_{h_u\tilde{\nu}}^2$. Of course, this last relation is not valid if the sneutrino vev's are not zero.

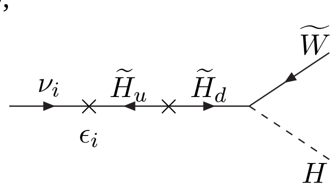
3. Change of Basis

Here we compare our results with previous publications. Consider first the MSSM case, *i.e.*, the scalars are not decoupled. In ref. [20], the neutral scalar loops contributing to the neutrino masses are studied. In their eq. (4.1) we see this contribution, and its main characteristic is that it vanishes in the decoupling limit. This is because the H and A contributions are negligible when m_H and m_A are very large, and because in the decoupling limit we have $\alpha = \beta - \pi/2$, canceling the contribution from the light Higgs boson h. As the authors explain above eq. (4.1), this result is written in the basis where the sneutrino vevs are zero, where the relation $M^2_{h_d\tilde{\nu}} = -\tan\beta M^2_{h_u\tilde{\nu}}$ holds. In this way, contributions from sneutrino mixing with up and down Higgs bosons can both be written proportional to B_iB_j, which appears as an overall factor. Other contributions from the same loops renormalize the tree level neutrino mass matrix, so they neglect them. The same result can be found in ref. [21] also written in the basis where the sneutrinos have zero vevs.

This is not contradictory with our result because we work in a basis where the sneutrino vevs are not zero. To better appreciate the differences, consider the case with only one leptonic superfield. The relevant terms in the superpotential are,

$$W \sim h_b\widehat{Q}\widehat{D}\widehat{H}_d - \mu\widehat{H}_d\widehat{H}_u + \epsilon_3\widehat{L}_3\widehat{H}_u \tag{B13}$$

where the bilinear ϵ_3 is the only term that violates R-Parity. We have omitted the terms we do not need, and we have considered only the third generation for simplicity. In this basis, it is well known that the tree-level neutrino mass matrix satisfy $M^{(0)}_{\nu\,ij} \sim \Lambda_i\Lambda_j$, where $\Lambda_i = \mu v_i + v_d\epsilon_i$ [6], generating the atmospheric mass. The Higgs boson loop that interests us here proceeds via the neutrino-wino-Higgs coupling, which can be understood in the following way,



$$\sim -\tfrac{\epsilon_i}{\mu}\tilde{g}c_\beta$$

where we use the two component formalism and the higgsino-wino-Higgs coupling is in eq. (1). This loop produces a one-loop contribution of the type $M^{(1)}_{\nu\,ij} \sim \epsilon_i\epsilon_j$, breaking the symmetry of the tree-level mass matrix, and generating a solar mass.

We look at these results now in a different basis. If we define $\widehat{L}'_3$ as the rotated leptonic superfield that has zero vev, then we need (see also [22])

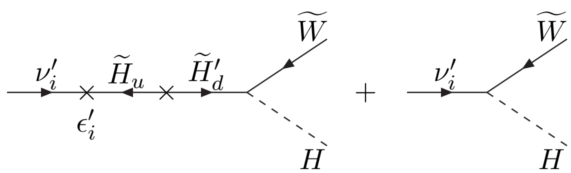
$$\widehat{H}'_d = \frac{v_d\widehat{H}_d + v_3\widehat{L}_3}{\sqrt{v_d^2 + v_3^2}}\,, \qquad \widehat{L}'_3 = \frac{-v_3\widehat{H}_d + v_d\widehat{L}_3}{\sqrt{v_d^2 + v_3^2}} \tag{B14}$$

Using the new fields, and defining $v_d'^2 = v_d^2 + v_3^2$, the superpotential is given by,

$$W \sim \frac{h_b v_d}{v_d'} \widehat{Q}\widehat{D}\widehat{H}_d' - \frac{h_b v_3}{v_d'} \widehat{Q}\widehat{D}\widehat{L}_3' - \frac{\mu v_d - \epsilon_3 v_3}{v_d'} \widehat{H}_d'\widehat{H}_u + \frac{\mu v_3 + \epsilon_3 v_d}{v_d'} \widehat{L}_3'\widehat{H}_u \tag{B15}$$

The new R-Parity violating parameter, which we could call $\epsilon_3' \equiv (\mu v_3 + \epsilon_3 v_d)/v_d'$ will contribute to the tree level neutrino mass $m_{\nu_3} \propto \epsilon_3'^2$, or to the neutrino mass matrix in the case of three generations, $M_\nu^{ij} \sim \epsilon_i'\epsilon_j'$. Since ϵ_i' parameters are proportional to Λ_i, this tree-level contribution is the same as in the previous basis.

In order to find the one-loop contribution in the rotated basis, we need the neutrino-wino-Higgs coupling. This time it is composed by two pieces,



$$\sim -\left(\frac{\epsilon_i'}{\mu} - \frac{v_i}{v_d}\right)\tilde{g} c_\beta$$

In this way, this vertex leads to the same result for the neutrino mass matrix as in the previous basis, that is $M_{\nu\, ij}^{(1)} \sim \epsilon_i\epsilon_j$. Note that the second contribution to the vertex comes from the Higgs-wino-higgsino term in eq. (1), and that we have neglected terms of second order in the estimation.

As can be seen in the superpotential in eq. (B15), in the rotated basis there is a trilinear R-Parity violating term of the type $\lambda' \sim h_b v_3/v_d'$, but it does not contribute to the neutrino mass matrix because squarks are too heavy. The pure BRpV model can then be understood in the basis with no sneutrino vevs by the existence of two contributions to the neutrino-wino-Higgs vertex that add up exactly equal to the vertex in the original basis.

Now we show how this is to be understood in the context of Split Supersymmetry. The low energy parameters a_i are related to the high energy parameters s_i through the matching condition at the scale $\widetilde{m}$ given in eq. (10). The parameters a_i need to be small due to neutrino physics constraints, typically of the order of ϵ/μ. At the scale $\widetilde{m}$ we have,

$$s_i = \frac{c_\beta M^2_{h_d\tilde{\nu}_i} + s_\beta M^2_{h_u\tilde{\nu}_i}}{M^2_{L_i} - m_h^2} \tag{B16}$$

where m_h is the light Higgs boson mass, which can be neglected in front of the sneutrino mass M_{L_i} in Split Supersymmetry. The Higgs-sneutrino mixings are given in eq. (B10). If we make the sneutrino vevs equal to zero, the mixings s_i would vanish as can be inferred from eqs. (B16), (B10) and (B12), and the BRpV contribution to the solar mass would be zero. On the contrary, if sneutrino vevs are not zero the minimization conditions in eq. (B12) imply that at the scale $\widetilde{m}$ we have $s_i \approx -s_\beta v_i/v_u$ and, therefore, $a_i \approx -v_i/v_d$. This is of the right order of magnitude for a correct solar mass. From the tadpole equation (B12) we see that the only requirement to have a non-zero sneutrino vev in Split Supersymmetry is $v_u B_\epsilon^i \sim v_i M^2_{L_i}$. In supergravity theories where $B^i_\epsilon = B_i\epsilon_i$ this implies an abnormally large value for B_i, disfavoring this kind of solution. An argument of this type can be found in ref. [8] where it is said that the large value for the sneutrino mass would make the Higgs-sneutrino mixing negligible. Here we argue that the origin for the term B^i_ϵ may be different, and we consider the former requirement as reasonable.

Finally, in ref. [23] the results valid for the basis where the sneutrino vevs are zero are generalized to an arbitrary basis. The difference between the basis-independent result in eq. (B20) of ref. [23] and the zero-sneutrino-vev basis result in eq. (4.4) of ref. [20] is the replacement of the term B_iB_j in [20] by a basis-independent form in [23] where $B_i \to |\vec{B}|(L^*L^T\delta_B)_i$, with δ_B and L defined in their Table I and eq. (3) respectively. Apart from the overall factor and sum, the body of the result remains unchanged. In particular, the contribution from the different Higgs bosons maintain their relative form such that in the decoupling limit the whole contribution vanishes. Since this relative form was obtained in the basis with zero sneutrino vevs (and clearly explained in [20]), using the relation $M^2_{h_d\tilde\nu} = -\tan\beta M^2_{h_u\tilde\nu}$, we argue that the result in eq. (B20) of ref. [23] is incomplete. The true relation between the up and down Higgs mixing with sneutrinos that should be used contains extra terms that vanish when the sneutrino vevs are zero. These missing terms are the ones that break the symmetry of the neutrino mass matrix, and can be found from the minimization condition in our eq. (B12), which coincides with eq. (3.8) in [20] in the limit where the slepton mass matrix is diagonal. We include these terms in our calculation, and that is why we are able to generate a solar mass.

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